Weak Radiative Decay $\Lambda_b \to \Lambda\ \gamma$ and Quark-Confined Effects in the Covariant Oscillator Quark Model

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Motivated by the observation of the decay $B \to K^* \gamma$ by the CLEO collaboration, we have systematically analyzed the weak radiative decay $\Lambda_b \to \Lambda\ \gamma$, evaluating the confined effects of quarks in the covariant oscillator quark model. This decay process receives both short distance (electromagnetic penguins at the one loop level) the long distance contributions in the quark level. The long distance contributions are analyzed using the vector meson dominance (VMD) method. The estimated branching ratio is found to be 0.23×10^{-5}.

§1. Introduction

Weak radiative (flavor changing neutral current) decays of hyperons have attracted the interest of physicists during the last three decades. In the Standard Model (SM) these processes are forbidden at the tree level and are strongly suppressed by the GIM mechanism. Hence they offer a unique possibility to test the CKM sector of the SM and possibly open a door to physics beyond it. Experimental data are now available for the light baryon sector, i.e., $\Sigma^+ \to p\ \gamma$, $\Lambda \to n\ \gamma$, $\Xi^- \to \Sigma^-\ \gamma$, $\Xi^0 \to \Sigma^0\ \gamma$ and $\Xi^0 \to \Lambda\ \gamma$,[1] which involve transitions of the type $s \to d\gamma$. Recently the rare decay $B \to K^*\ \gamma$ has been observed by the CLEO collaboration,[2] which is dominated by the quark level process $b \to s\gamma$. Therefore, one may expect that there is also some possibility for the rare decay of heavy baryons. For the heavy hadron decay processes of this type, it is considered that the confined effects of quarks play generally an important role, since there is a large difference between the initial and final hadron masses. In a preceding paper[3] we investigated the rare $B \to K^*\gamma$ process, evaluating the confined effects in the framework of the covariant oscillator quark model (COQM).[4]

In this paper we study the weak radiative decay $\Lambda_b \to \Lambda\ \gamma$, which results from the interplay of electroweak and gluonic interactions. At the quark level, there are two essential mechanisms responsible for the weak radiative decays $(b \to s\gamma)$: the short distance (SD) electroweak penguin and the long distance (LD) contributions. Recently, an investigation[5] of long distance contributions to these decays was made using vector meson dominance (VMD) at the quark level as $b \to s[\psi]$ followed by the conversion $[\psi] \to \gamma$. We estimate the contributions to the decay $\Lambda_b \to \Lambda\ \gamma$ arising

from both sources, evaluating the hadronic matrix elements in COQM. [4]

One of the most important motivations for the COQM is to covariantly describe the center of mass motion of hadrons, preserving the considerable successes of the nonrelativistic quark model on the static properties of hadrons. A key in COQM for doing this is to treat the square masses of hadrons as opposed to the mass itself, as done in conventional approaches. This makes the covariant treatment simple. The COQM has been applied to various problems [6] with satisfactory results. Recently, Ishida et al. [7],[8] have studied the weak decays of heavy hadrons using this model and derived the same relations of weak form factors for heavy-to-heavy transitions as in HQET. [9] In addition, the COQM is also applicable to heavy-to-light transitions. As a consequence, this model does incorporate the features of heavy quark symmetry and can be used to compute the form factors for heavy-to-light transitions as well, which is beyond the scope of HQET. Actually, in previous papers we made analyses of the spectra of exclusive semi-leptonic [8] decays of B-mesons, of non-leptonic decays of B mesons, [10] and of hadronic weak decays of Λ_b baryons [11] along this line of reasoning, leading to encouraging results.

The paper is organised as follows. In §2 we present the methodology necessary for our analysis. The short distance and long distance effects are discussed in §2.1 and 2.2 and evaluation of the hadronic matrix elements are given in §2.3. Section 3 contains our results and discussion.

§2.   Methodology

The general amplitude of baryon weak radiative decay is given by

$$\mathcal{M}(B_i \to B_f \ \gamma) = i\bar{u}_f(a + b\gamma_5)\sigma_{\mu\nu}\epsilon^\mu k^\nu u_i \qquad (2\cdot1)$$

where u_i (u_f) is the Dirac spinor of the initial (final) baryon, ϵ^μ (k^ν) denotes the polarization (momentum) vector of the photon, and a and b are parity-conserving and parity-violating amplitudes, respectively. The corresponding decay rate is given as

$$\Gamma(B_i \to B_f \ \gamma) = \frac{1}{8\pi}\left(\frac{M_i^2 - M_f^2}{M_i}\right)^3 (|a|^2 + |b|^2) , \qquad (2\cdot2)$$

and the asymmetry parameter is given as

$$\alpha = \frac{2\mathrm{Re}\ (a^*b)}{(|a|^2 + |b|^2)} \qquad (2\cdot3)$$

2.1.   *Short distance contribution*

The effective Hamiltonian [12] for the short distance $b \to s$ transition including the QCD correction is given by

$$\mathcal{H}_{\mathrm{eff}}^{SD}(b \to s\gamma) = -\frac{G_F}{\sqrt{2}}\frac{e}{16\pi^2}F_2 V_{tb}V_{ts}^* F_{\mu\nu}\ [m_b\bar{s}\ \sigma^{\mu\nu}\ (1 + \gamma_5)\ b + m_s\bar{s}\ \sigma^{\mu\nu}\ (1 - \gamma_5)\ b] ,$$

$$(2\cdot4)$$

where $F_{\mu\nu}$ is the electromagnetic field strength tensor, V_{ij} are the CKM matrix elements, $F_2 \approx F_2(x_t) - F_2(x_c) \approx F_2(x_t)$ with $x_i = m_i^2/M_W^2$ and

$$F_2(x) = \rho^{-16/23} \left\{ \bar{F}_2(x) + \frac{116}{27} \left[\frac{1}{5}(\rho^{10/23} - 1) + \frac{1}{14}(\rho^{28/23} - 1) \right] \right\} \tag{2·5}$$

with

$$\bar{F}_2(x) = \frac{(8x^2 + 5x - 7)x}{12(x-1)^3} - \frac{(3x-2)x^2}{2(x-1)^4}\ln x \,, \tag{2·6}$$

and

$$\rho = \frac{\alpha_s(m_b^2)}{\alpha_s(M_W^2)} = 1 + \frac{23}{12\pi}\alpha_s(m_b^2)\,\ln\left(\frac{M_W^2}{m_b^2}\right) \,. \tag{2·7}$$

Numerically $F_2(x_t) = 0.65$ for $\Lambda_{\mathrm{QCD}} = 200$ MeV and $m_t = 174$ GeV. Furthermore, using the unitarity of the CKM matrix, i.e., $V_{tb}V_{ts}^* = -V_{cs}^*V_{cb} - V_{us}^*V_{ub}$ since $V_{us}^*V_{ub} << V_{cs}^*V_{cb}$, we obtain the matrix element arising from the short distance effects as

$$\langle \Lambda\gamma|\mathcal{H}_{\mathrm{eff}}^{SD}|\Lambda_b\rangle = i\frac{G_F}{\sqrt{2}}\frac{e}{8\pi^2}F_2 V_{cb}V_{cs}^*\epsilon_\mu k_\nu\langle\Lambda|\bar{s}[m_b\sigma^{\mu\nu}(1+\gamma_5) + m_s\sigma^{\mu\nu}(1-\gamma_5)]b|\Lambda_b\rangle \,. \tag{2·8}$$

2.2. *Long distance contribution*

The long distance contributions were recently estimated using the vector meson dominance method in Ref. 5). At the quark level, this assumes the dominance of the process $b \to s[\sum_i \psi_i] \to s\,\gamma$, where all the $\bar{c}c$, $J = 1$ excited as well as ground charmonium states are taken into acount as ψ_i. The relevant part of the effective Hamiltonian describing the process is given by

$$\mathcal{H}_{\mathrm{eff}} = \frac{G_F}{\sqrt{2}}\,V_{cb}\,V_{cs}^*\,[C_1(\mu)\,O_1 + C_2(\mu)\,O_2], \tag{2·9}$$

with

$$O_1 = (\bar{s}c)^\mu\,(\bar{c}b)_\mu \qquad \text{and} \qquad O_2 = (\bar{s}b)^\mu(\bar{c}c)_\mu \,, \tag{2·10}$$

where the quark current $(\bar{q}_i q_j)_\mu = \bar{q}_i\gamma_\mu(1-\gamma_5)q_j$ denotes the usual $(V-A)$ current.

Using factorization, we obtain the inclusive decay amplitude for the process $b \to s\psi$ as

$$\mathcal{M}(b \to s\psi(k_1,\epsilon_1)) = \frac{G_F}{\sqrt{2}}V_{cs}^*V_{cb}a_2(\mu)f_\psi(m_\psi^2)m_\psi\bar{s}\gamma^\mu(1-\gamma_5)b\epsilon_{1\mu} \,, \tag{2·11}$$

where $a_2(\mu) = C_2(\mu) + C_1(\mu)/N_c$, and $N_c = 3$ is the number of colors. k_1 and ϵ_1 are the momentum and polarization vector of the vector meson ψ. In the above equation we have used the matrix element

$$\langle\psi(k_1,\epsilon_1)|(\bar{c}c)^\mu|0\rangle = f_\psi(m_\psi^2)\,m_\psi\epsilon_1^\mu. \tag{2·12}$$

We have used the value a_2^{eff}, which is determined experimentally from the world average branching ratio of $\bar{B} \to K^* \psi$ [13] as

$$a_2^{\text{eff}} = 0.23 \ . \tag{2.13}$$

Now we wish to replace the ψ meson with the photon and construct a gauge invariant amplitude. This can be done by eliminating the longitudinal component of the ψ meson so that ϵ_1^μ is changed to the polarization vector of the photon ϵ^μ. For this purpose we use the procedure of Golowich and Pakvasa. [14] Now, using the equation of motion for the b quark, i.e. $\not{p}b = m_b b$, and momentum conservation, $p = p' + k_1$, we obtain

$$\bar{s}\gamma_\mu(1 - \gamma_5)b = \frac{1}{m_b}[\bar{s}\gamma_\mu \not{p}'(1 + \gamma_5)b + \bar{s}\gamma_\mu \not{k}_1(1 + \gamma_5)b], \tag{2.14}$$

where p and p' are the momenta of the b and s quarks respectively. The contribution due to the first term in Eq. (2.14) is neglected, since $m_s \ll m_b$ and $p'^\mu \epsilon_{1\mu}^T = 0$, which follows from $\epsilon_{1\mu}^T p^\mu = 0$ in the rest frame of the b quark and the transversality condition $\epsilon_{1\mu}^T k_1^\mu = 0$, where $\epsilon_{1\mu}^T$ is the transverse polarization vector of the ψ meson. The second term can be written as

$$\frac{1}{m_b}\bar{s}\gamma_\mu \not{k}_1(1 + \gamma_5)b = \frac{1}{m_b}\{\bar{s}(1 + \gamma_5)k_{1\mu}b - i\bar{s}\sigma_{\mu\nu}k_1^\nu(1 + \gamma_5)b\} \ . \tag{2.15}$$

In Eq. (2.15) only the $\sigma_{\mu\nu}$ term couples to the transverse component of ψ, and we obtain the corresponding amplitude as

$$\mathcal{M}(b \to s\psi^T) = -\frac{G_F}{\sqrt{2}}V_{cb}V_{cs}^* a_2(\mu)f_\psi(m_\psi^2)\frac{m_\psi}{m_b}\bar{s}\sigma_{\mu\nu}(1 + \gamma_5)b\epsilon_1^{T\mu}k_1^\nu \ . \tag{2.16}$$

For the $\psi^T \to \gamma$ conversion following the VMD mechanism, we have

$$\langle 0|J_\mu|\psi^T(k_1, \epsilon_1^T)\rangle = eQ_c f_\psi(0)m_\psi\epsilon_{1\mu}^T, \tag{2.17}$$

where $Q_c = 2/3$ and $f_\psi(0)$ is the coupling at $k_1^2 = 0$. Using the intermediate propagator of the ψ meson at $k_1^2 = 0$, we get

$$\mathcal{M}(b \to s\psi^T \to s\gamma) = i\frac{G_F}{\sqrt{2}} V_{cb}V_{cs}^* a_2 f_\psi^2(0) \frac{eQ_c}{m_b} \bar{s} \sigma_{\mu\nu} (1 + \gamma_5) b \epsilon^\mu k^\nu \ . \tag{2.18}$$

It should be noted that the coupling structure is the same as that due to the short distance electromagnetic penguin operator. The expression for the amplitude Eq. (2.18) can be completed by summing over all the $c\bar{c}$ resonant states $\psi(1S)$, $\psi(2S)$, $\psi(3770)$, $\psi(4040)$, $\psi(4160)$ and $\psi(4415)$:

$$\mathcal{M}(b \to s\psi^T \to s\gamma) = i\frac{G_F}{\sqrt{2}} V_{cb}V_{cs}^* a_2 \kappa \sum_i f_{\psi_i}^2(m_{\psi_i}^2) \frac{eQ_c}{m_b} \bar{s} \sigma_{\mu\nu} (1 + \gamma_5) b \epsilon^\mu k^\nu \ . \tag{2.19}$$

The various decay couplings $f_{\psi_i} = f_{\psi_i}(m_{\psi_i}^2)$ are calculated using

$$f_{\psi_i}^2 = \Gamma(\psi_i \to e^+ e^-)\frac{3m_{\psi_i}}{Q_c^2 4\pi\alpha^2}, \tag{2.20}$$

Table I. Values of the vector meson decay constants

ψ_i	$\psi(1S)$	$\psi(2S)$	$\psi(3770)$	$\psi(4040)$	$\psi(4160)$	$\psi(4415)$
f_{ψ_i} (GeV)	0.405	0.282	0.099	0.175	0.180	0.145

which are given in Table I. To extrapolate the coupling $f_{\psi_i}(k_1^2 = m_\psi^2)$ to $f_{\psi_i}(0)$, we use the suppression factor [5]

$$\kappa = \frac{f_{\psi(1S)}^2(0)}{f_{\psi(1S)}^2(m_\psi^2)} = 0.12 \; , \tag{2.21}$$

obtained from the data on the photoproduction of the ψ. This is taken to be universal for all resonances. Now, we use Eq. (2·19) to find the matrix element for $\Lambda_b \to \Lambda\gamma$ through the $b \to s\psi^T \to s\gamma$ transition at the quark level, which is given as

$$\langle \Lambda \, \gamma | \mathcal{H}_{\text{eff}}^{\bar{c}c} | \Lambda_b \rangle = i\frac{G_f}{\sqrt{2}} V_{cb} V_{cs}^* a_2 e Q_c \kappa \epsilon^\mu k^\nu \sum_i \frac{f_{\psi_i}^2(m_{\psi_i}^2)}{m_b} \langle \Lambda | \bar{s}\sigma_{\mu\nu}(1+\gamma_5)b | \Lambda_b \rangle \; . \tag{2.22}$$

2.3. *Evaluation of the hadronic form factors*

In this section we evaluate the hadronic matrix elements present in the expression for the decay amplitudes Eqs. (2·8) and (2·22). However, these hadronic matrix elements involve the tensor currents, and there seems to be no known method to evaluate them. On the other hand, these elements are easily evaluated [11), 15)] in COQM by taking the overlapping of the initial and final wave functions. For Λ_Q-type baryons, udQ, the Bargmann-Wigner spinor function of the third constituent quark Q changes into the spinor wave functions of the total baryon wave functions. Thus, the tensor current is calculated as [15]

$$\langle \Lambda | \bar{s}i\sigma_{\mu\nu}(1 \pm \gamma_5)b | \Lambda_b \rangle = I_{ud}^{sb}(w)\bar{u}_f \frac{w+1}{2} i\sigma_{\mu\nu}(1 \pm \gamma_5)u_i \; . \tag{2.23}$$

The single form factor function $I_{ud}^{sb}(w = -v \cdot v')$ denotes the overlapping of the initial and final space-time wave function. It describes the confinement effects of the quarks and is given by

$$I_{ud}^{sb}(w) = \frac{1}{w} \frac{4\beta_\lambda\beta_\lambda'}{\beta_\lambda + \beta_\lambda'} \frac{1}{\sqrt{C(w)}} \exp(-G(w)) \; , \tag{2.24}$$

where

$$C(w) = (\beta_\lambda - \beta_\lambda')^2 + 4\beta_\lambda\beta_\lambda' \, w^2 \; , \tag{2.25}$$

and

$$G(w) = \frac{4m_q^2(\beta_\lambda + \beta_\lambda') \, w(w-1)}{(\beta_\lambda - \beta_\lambda')^2 + 4\beta_\lambda\beta_\lambda' \, w^2} \; . \tag{2.26}$$

Here β_λ and β_λ' are the oscillator strength of the initial and final baryon oscillator wave functions. They are given in terms of the quark masses as

$$\beta_\lambda = \sqrt{\frac{m_q m_b K}{2m_q + m_b}} \qquad \text{and} \qquad \beta_\lambda' = \sqrt{\frac{m_q m_s K}{2m_q + m_s}} \; . \tag{2.27}$$

In the above equation, m_q, m_s and m_b denote the masses of the u/d, s and b quarks, and K is the universal spring constant for all hadronic systems with the value $K = 0.106 \text{ GeV}^3$. [16)]

§3. Results and conclusion

Thus with Eqs. (2·1), (2·8), (2·22) and (2·23) we find the parity conserving and parity violating ampliudes a and b as

$$a = \frac{G_F}{\sqrt{2}}eV_{cb}V_{cs}^*\frac{w+1}{2}I_{ud}^{sb}(w)\left[\frac{F_2}{8\pi^2}(m_b+m_s) + \frac{2}{3}a_2\kappa\sum_i\frac{f_{\psi_i}^2(m_\psi^2)}{m_b}\right] \tag{3·1}$$

and

$$b = \frac{G_F}{\sqrt{2}}eV_{cb}V_{cs}^*\frac{w+1}{2}I_{ud}^{sb}(w)\left[\frac{F_2}{8\pi^2}(m_b-m_s) + \frac{2}{3}a_2\kappa\sum_i\frac{f_{\psi_i}^2(m_\psi^2)}{m_b}\right] \tag{3·2}$$

To estimate the numerical result we have used the following values. The quark masses used are $m_q = 0.4$ GeV, $m_s = 0.51$ GeV and $m_b = 5$ GeV, which were determined through the analysis of mass spectra. [16)] The particle masses and the lifetime of the Λ_b baryon are taken from Ref. 1). The values of the CKM matrix elements used are $V_{cb} = 0.0395$ and $V_{cs} = 1.04$. With these values, we obtain the branching ratio for the decay $\Lambda_b \to \Lambda\gamma$ to be

$$Br(\Lambda_b \to \Lambda\gamma) = 0.23 \times 10^{-5} . \tag{3·3}$$

The main contribution comes from the SD amplitude. By taking only the SD amplitude into account, the branching ratio is predicted to be 0.215×10^{-5}.[*)] In obtaining the above branching ratio, Eq. (3·3), the value of the form factor function was $I_{ud}^{sb}(\omega = 2.64) = 0.0395$. This extremely small value compared with $I = 1$ (corresponding to the "free-quark decay") indicates that the quark-confined effects largely reduce the decay rate of $\Lambda_b \to \Lambda\,\gamma$.

The asymmetry parameter is predicted as

$$\alpha = 0.98 . \tag{3·4}$$

Our predicted branching ratio lies slightly below the theoretical value [17)] $(1 \pm 0.5) \times 10^{-5}$, while our predicted value of α is consistent with the theoretical value [17)] 0.9. Additional experimental data would greatly help for a better understanding of the weak radiative decays of heavy baryons, which can serve as a signal of new physics beyond the Standard Model.

This decay process was previously studied by Cheng et al. [18)] They took into account the short distance contribution only and estimated the branching ratio following two different approaches. In the first method they treated both the b and s quarks as heavy and included a correction of order $1/m_s$. However, the $1/m_s$

[*)] The branching ratio due only to the LD contribution is 1.935×10^{-9}, which is very small in comparison with the value due only to the SD contribution.

correction to the $\Lambda_b \to \Lambda\gamma$ amplitude is about 50% for $m_s = 510$ MeV which is quite sizable. Hence it is important to include higher order $1/m_s$ corrections. In the second method they treated the b quark as heavy and used the MIT bag model to calculate the form factors at zero recoil and then extrapolate them to $q^2 = 0$ by assuming a dipole q^2 dependence of the form factor. In this paper we have considered both the short and long distance contributions and taken into account the confined effects by using the COQM, which has proven to be very successful for the phenomenology of both the heavy and light hadrons. It should be noted that in this model we do not need to extrapolate the form factor to the particular point of interest, as the model gives the full spectrum for the hadronic form factor functions. Thus our model provides a better theoretical understanding of the rare decays of heavy baryons.

Acknowledgements

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